Supersymmetry breaking and loop corrections at the end of inflation

R. Jeannerot[1] and J. Lesgourgues[2]

[1]International Centre for Theoretical Physics (ICTP), P.O. Box 586, Strada Costiera 11, 34014 Trieste, Italy;
[2] SISSA/ISAS, Via Beirut 2-4, 34014 Trieste, Italy

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We show that quantum corrections to the effective potential in supersymmetric hybrid inflation can be calculated all the way from the inflationary period - when the Universe is dominated by a false vacuum energy density - till the fields settle down to the global supersymmetric minimum of the potential. These are crucial for getting a continuous description of the evolution of the fields.

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I. INTRODUCTION

Inflation solves many of the outstanding problems of the standard cosmology [1]. Among others, it provides a mechanism for the generation of primordial cosmological perturbations, which are responsible for the observed temperature anisotropies in the cosmic microwave background (CMB), and the large-scale structure (LSS) in our Universe. Future experiments, such as MAP* and Planck† (resp. 2dF‡ and SDSS§), will measure with great precision the power spectrum of the CMB (resp. LSS), and draw sharp constraints on the potential of the inflaton, the scalar field driving inflation.

Successful implementation of the inflationary picture requires a long enough era of accelerated expansion on one hand, and a correct order of magnitude for primordial perturbations on the other hand. In the most simple versions of single-field inflation, the corresponding constraints on the inflaton potential are unrealistic from a particle physics point of view. Indeed, coupling parameters must be fine-tuned to very small values, while the inflaton must be of the order of the Planck mass during inflation. These problems can be avoided in the so-called hybrid models [2] (see also Ref. [3]), in which the inflaton couples with some other(s) scalar field(s). Hybrid inflation arises naturally in supersymmetric theories** [3,5–7]. Supersymmetry provides the flatness of the scalar potential required for inflation. The inflaton field, which is usually a scalar singlet, couples to Higgs superfield(s) charged under some gauge group G. At the end of inflation, the Higgs fields acquire a non-vanishing vacuum expectation value (VEV); G is spontaneously broken. Such scenarios arise naturally in supersymmetric grand unified theories (SUSY GUT) [5,8]. The non-zero vacuum energy density during inflation can either be due to the VEV of a F-term [3,5] or from that of a D-term [3,6,7].

Models of either types share the following features. The scalar potential has two minima; one local minimum for values of the inflaton field $|S|$ greater than some critical value s_c with the Higgs fields at zero, and one global supersymmetric minimum at $S = 0$ with non-zero Higgs VEVs. The fields are usually assumed to have chaotic initial conditions [9], with an initial value $|S| \gg s_c$ for the inflaton. The Higgs fields rapidly settle down to the local minimum of the potential††; then, the universe is dominated by a non-vanishing vacuum energy density, and supersymmetry is broken. This in turn leads to quantum corrections to the potential which lift its complete flatness [5,11]. The slow-roll conditions are satisfied and inflation takes place until $|S| = s_c$ or slighlty before, depending on the model. When $|S|$ falls below s_c, the Higgs fields start to acquire non-vanishing VEVs. All fields then oscillate until they stabilise at the global supersymmetric minimum. These oscillations are crucial for understanding the process of reheating and particle production in the early universe.

The important point is that, as long as the fields are not settled down at the global minimum, supersymmetry remains broken. When the fields oscillate, the system only goes ponctually through supersymmetric configurations. The breaking of supersymmetry is best seen by looking at the mass spectrum; the bosonic and fermionic masses are non degenerate. This has important consequences. It implies that loop corrections to the effective potential are non-zero not only during inflation, but also during all the oscillatory regime. The corrections are crucial for getting a continuous description of the evolution of the fields. They will be useful for the simulation of preheating, for the calculation of the number density of cosmic strings, for the study of leptogenesis at the end of

*http://map.gsfc.nasa.gov

†http://astro.estec.esa.nl/SA-general/Projects/Planck

‡http://meteor.anu.edu.au/ colless/2dF

§http://www.astro.princeton.edu/BBOOK

**For a review on inflation in supersymmetric theories see Ref. [4].

††The problem of initial conditions in inflation is actually not trivial. We shall not discuss this problem here. See Refs. [10] for possible solutions.

inflation [12,13], and for the derivation of the primordial spectrum during intermediate stages in supersymmetric multiple inflationary models [14,15].

In this Letter, we calculate the one-loop corrections to the potential along the inflaton direction. They are the most important and affect the dynamics of the inflaton field, which in turn affects the dynamics of the Higgs fields. They can be calculated by applying the Coleman-Weinberg formula [16]:

$$\Delta V = \frac{1}{64\pi^2}\sum_i (-1)^F m_i^4 \ln(m_i^2/\Lambda^2), \tag{1}$$

where $(-1)^F$ shows that bosons and fermions make opposite contributions; it is $+1$ for the bosonic degrees of freedom and -1 for the fermionic ones. The sum runs over each degree of freedom i with mass m_i and Λ is a renormalization scale. We thus determine the particle spectrum for each value of the inflaton field $|S|$ and values of the other fields which minimize the potential for this $|S|$. We consider the standard models of F and D-term inflation. The particle spectrum is found to be very rich and interesting. During inflation, the non-zero quantum corrections are due to a boson-fermion mass splitting in the Higgs sector. When $|S|$ falls below s_c, since the Higgs VEVs are non zero, there is also a mass splitting between gauge and gaugino fields.

II. F-TERM INFLATION

The simplest superpotential which leads to F-term inflation is given by $W = \alpha S\overline{\Phi}\Phi - \mu^2 S$, where S is a scalar singlet and (Φ, $\overline{\Phi}$) are Higgs superfields in complex conjugate representations of some gauge group G [3,5]. α and μ are two constants which are taken to be positive and $\frac{\mu}{\sqrt{\alpha}}$ sets the G symmetry breaking scale. This superpotential is consistent with a continuous R symmetry under which the fields transform as $S \to e^{i\gamma}S$, $\Phi \to e^{i\gamma}\Phi$, $\overline{\Phi} \to e^{-i\gamma}\overline{\Phi}$ and $W \to e^{i\gamma}W$. It is is often used in SUSY GUT model building. The scalar potential reads:

$$\begin{aligned} V = {} & \alpha^2|S|^2(|\overline{\Phi}|^2 + |\Phi|^2) + |\alpha\overline{\Phi}\Phi - \mu^2|^2 \\ & + \frac{g^2}{2}(|\overline{\Phi}|^2 - |\Phi|^2)^2, \end{aligned} \tag{2}$$

where we have kept the same notation for the superfields and their bosonic components. There is a flat direction with degenerate local minima $|S| \equiv s \geq s_c = \frac{\mu}{\sqrt{\alpha}}$, $\overline{\Phi} = \Phi = 0$, for which $V = \mu^4$, and a global supersymmetric minimum at $S = 0$, $|\Phi| = |\overline{\Phi}| = \frac{\mu}{\sqrt{\alpha}}$, $\arg(\Phi)+\arg(\overline{\Phi}) = 0$, in which the G symmetry is spontaneously broken.

We assume that the problem of initial conditions has been solved, and we investigate the behavior of the system already settled in the local minimum of the potential. The universe is dominated by the vacuum energy density $V = \mu^4$ and supersymmetry is broken. The bosonic and fermionic masses are thus non degenerate. The mass splitting happens in the Φ and $\overline{\Phi}$ sector. Explicitly, there are two complex scalars with masses squared $m_1^2 = \alpha^2 s^2 + \mu^2\alpha$ and $m_2^2 = \alpha^2 s^2 - \mu^2\alpha$ (the mass eigenstates are linear combinaisons of the Φ and $\overline{\Phi}$ fields), and two Weyl fermions with masses $m^2 = \alpha^2 s^2$. This spectrum gives rise to quantum corrections to the effective potential which can be calculated from Eq.(1) and lift the the complete flatness of the s direction. When $s \gg s_c$, they have a well-know asymptotic form [5]:

$$\Delta V = \frac{\alpha^2\mu^4}{16\pi^2}\left(\ln\frac{\alpha^2 s^2}{\Lambda^2} + \frac{3}{2}\right). \tag{3}$$

Therefore, the S field can roll down the potential. The slow roll conditions are satisfied and inflation takes place. When s falls below s_c, Φ and $\overline{\Phi}$ are destabilized, and all fields start to oscillate.

During inflation, the Higgs fields Φ and $\overline{\Phi}$ have zero VEVs and since the inflaton S is assumed to be a gauge singlet, gauge and gauginos have zero masses; the only contribution to ΔV comes from the mass splitting in the Φ and $\overline{\Phi}$ sector. Now when s falls below s_c, the VEVs of Φ and $\overline{\Phi}$ start to be non-zero. Thus the corresponding gauge and gaugino fields, as well as the S field, also acquire non-zero mass. The mass splitting then happens both in the Higgs and in the gauge sectors.

From now on, we shall assume that the VEVs of Φ and $\overline{\Phi}$ only break a U(1) gauge symmetry, and that the representation of Φ is complex one-dimensional. For arbitrary n-dimensional complex conjugate representations which break a gauge group G down to a subgroup H of G, when the Higgs VEVs are non-zero, there are $k = \dim(G) - \dim(H)$ massive gauge fields and $4n - k + 2$ massive real scalar fields. For any value of s, the potential is minimised along the D-flat direction $|\overline{\Phi}| = |\Phi| = \hat{\phi}$, and for $\arg(\Phi) = -\arg(\overline{\Phi}) = \theta$. Therefore, we expand the fields as follows:

$$\Phi = \hat{\phi}e^{i\theta} + \phi_1, \qquad \overline{\Phi} = \hat{\phi}e^{-i\theta} + \phi_2, \tag{4}$$

where ϕ_1 and ϕ_2 are complex fields which represent the quantum fluctuations of the Φ and $\overline{\Phi}$ fields. $\hat{\phi} = 0$ for $s \geq s_c$ and $\hat{\phi} = \sqrt{\frac{\mu^2 - \alpha s^2}{\alpha}}$ for $s \leq s_c$. When s falls below s_c, we find the following spectrum. There is a complex scalar field with squared mass $m_S^2 = 2\alpha^2\hat{\phi}^2$. The Higgs mechanism gives rise to three real scalars with masses squared $m_1^2 = 2\alpha^2\hat{\phi}^2$, $m_2^2 = 2\alpha\mu^2$ and $m_3^2 = 2\alpha^2 s^2 + 4g^2\hat{\phi}^2$. The corresponding Higgs mass eigenstates are $\Re(\phi_1)+\Re(\phi_2)$, $\Im(\phi_1)+\Im(\phi_2)$ and $\Re(\phi_1)-\Re(\phi_2)$. The field $\Im(\phi_1) - \Im(\phi_2)$ is absorbed by the gauge field which is now massive with mass squared $m_A^2 = 4g^2\hat{\phi}^2$.

The fermionic spectrum can be derived from the following parts of the Lagrangian:

$$\begin{aligned} \mathcal{L}_Y &= \alpha(S\psi_1\psi_2 + \Phi\psi_S\psi_2 + \overline{\Phi}\psi_1\psi_S), \\ \mathcal{L}_g &= -i\sqrt{2}g(\tilde{\Lambda}\psi_2\overline{\Phi}^* - \tilde{\Lambda}\psi_1\Phi^*) + \text{h.c.}, \end{aligned} \tag{5}$$

Field	d.o.f.	squared mass
$\frac{\phi_1^*+\phi_2}{\sqrt{2}}$	2	$\alpha^2 s^2 + \mu^2\alpha$
$\frac{-\phi_1^*+\phi_2}{\sqrt{2}}$	2	$\alpha^2 s^2 - \mu^2\alpha$
2 Weyl fermions	2×2	$\alpha^2 s^2$

TABLE I. Particle spectrum during F-term inflation and numbers of degrees of freedom (d.o.f.), when $s = s \geq s_c$.

where ψ_1, ψ_2 and ψ_S are the fermionic components of the Higgs and inflaton superfields. $\tilde{\Lambda}$ is the gaugino. After diagonalizing the fermion mass matrix, we find that there are four Weyl fermions with masses:

$$m^2_{\psi_1^\pm} = 2\alpha^2\hat{\phi}^2 + \frac{\alpha^2 s^2}{2} \pm \frac{1}{2}\alpha^2 s\sqrt{8\hat{\phi}^2 + s^2},$$
$$m^2_{\psi_2^\pm} = 4g^2\hat{\phi}^2 + \frac{\alpha^2 s^2}{2} \pm \frac{1}{2}\alpha s\sqrt{16g^2\hat{\phi}^2 + \alpha^2 s^2}. \qquad (6)$$

The mass eigenstates are, respectively, linear combinations of the higgsinos and the inflatino (the fermionic component of the inflaton superfield), and linear combinations of the higgsinos and the gaugino. We summarize our results in Tables I and II. The one-loop corrected effective potential is $V + \Delta V(s)$, where $\Delta V(s)$ is given by Eq. (1). We check that the supertrace $\mathrm{Str}M^2 \equiv \sum_i (-1)^F m_i^2$ vanishes at all times [17], and that the corrections are continuous at $s = s_c$. The exact effective potential should be a smooth function of s, and independent of Λ. In the one-loop approximation, Λ must be chosen so that the contribution of higher order terms can be neglected. This is generally achieved with $\Lambda^2 \sim \alpha^2 s_c^2 = \mu^2\alpha$ [18]. Here, by imposing the continuity of the potential derivative at $s = s_c$, we find:

$$\Lambda^2 = e^\epsilon \alpha^2 s_c^2, \qquad \epsilon \equiv \frac{1}{2} - \frac{g^2 \ln 2}{\alpha^2 + g^2}. \qquad (7)$$

When $g \leq \alpha$, the shape of $\Delta V(s)$ is given in Fig. 1, and $|\Delta V(s)| \ll V$. When $g > \alpha$, a higher mass scale $g^2\hat{\phi}^2 > \alpha^2 s_c^2$ appears after symmetry breaking; thus, the above choice of Λ^2 is not valid anymore, and we find that the one-loop corrections exceed the tree-level potential; so, in the one-loop approximation, it is impossible to find an expression for $\Delta V(s)$ valid around s_c.

III. D-TERM INFLATION

A generic toy-model of D-term inflation, proposed in Refs. [7], involves three complex fields: a gauge singlet S, and two fields Φ_+ and Φ_- with charges $+1$ and -1 under a $U(1)$ gauge symmetry. The superpotential is $W = \lambda S \Phi_+ \Phi_-$. This choice can be justified by a set of continuous R-symmetries or discrete symmetries. The scalar potential is:

Field	d.o.f.	squared mass
S	2	$2\alpha^2\hat{\phi}^2$
$\Re(\phi_1) + \Re(\phi_2)$	1	$2\alpha^2\hat{\phi}^2$
$\Im(\phi_1) + \Im(\phi_2)$	1	$2\alpha\mu^2$
$\Re(\phi_1) - \Re(\phi_2)$	1	$2\alpha^2 s^2 + 4g^2\hat{\phi}^2$
A	3	$4g^2\hat{\phi}^2$
2 Weyl fermions	2×2	$2\alpha^2\hat{\phi}^2 + \frac{\alpha^2 s^2}{2} \pm \frac{\alpha^2 s}{2}\sqrt{8\hat{\phi}^2 + s^2}$
2 Weyl fermions	2×2	$4g^2\hat{\phi}^2 + \frac{\alpha^2 s^2}{2} \pm \frac{\alpha s}{2}\sqrt{16g^2\hat{\phi}^2 + \alpha^2 s^2}$

TABLE II. Particle spectrum along the inflaton direction at the end of F-term inflation, when $\hat{\phi} = \sqrt{\frac{\mu^2 - \alpha s^2}{\alpha}}$ and $s = s \leq s_c$.

$$V = \lambda^2|S|^2(|\Phi_+|^2 + |\Phi_-|^2) + \lambda^2|\Phi_+|^2|\Phi_-|^2$$
$$+\frac{g^2}{2}(|\Phi_+|^2 - |\Phi_-|^2 + \xi)^2, \qquad (8)$$

where ξ is a Fayet-Illiopoulos term. We suppose that $\xi > 0$ (if it is not the case, the role of Φ_+ and Φ_- are just inverted). There is a true supersymmetric vacuum at $S = \Phi_+ = 0$, $|\Phi_-| = \sqrt{\xi}$, and a valley of local minima $|S| \equiv s > s_c = g\sqrt{\xi}/\lambda$, $\Phi_+ = \Phi_- = 0$, in which the tree-level potential is flat, $V = g^2\xi^2/2$, and the dynamics of s is only driven by quantum corrections. In the false vacuum, (Φ_+, Φ_-) have squared masses $\lambda^2 s^2 \pm g^2\xi$, while their fermionic superpartners (ψ_+, ψ_-) combine to form a Dirac spinor with mass λs (see Table III). The loop corrections are given by Eq. (1), and when $s \gg s_c$ they reduce to:

$$\Delta V = \frac{g^4\xi^2}{16\pi^2}\left(\ln\frac{\lambda^2 s^2}{\Lambda^2} + \frac{3}{2}\right). \qquad (9)$$

When s falls below s_c, Φ_- acquires a non-vanishing VEV $\hat{\phi}e^{i\theta}$, with $\hat{\phi} = \sqrt{\xi - \frac{\lambda^2}{g^2}s^2}$, which breaks the $U(1)$ gauge symmetry. The mass splitting then also happens in the gauge sector. Without loss of generality, we assume that $\theta = 0$, and expand the Higgs field as $\Phi_- = \hat{\phi} + \phi_1$. The real field $\sqrt{2}\Re(\phi_1)$ has a squared mass $2g^2\hat{\phi}^2$, while the Goldstone boson $\sqrt{2}\Im(\phi_1)$ is eaten up by the gauge boson, which becomes massive with $m_A^2 = 2g^2\hat{\phi}^2$. The masses for S and Φ_+ are given in Table IV. The fermionic spectrum can be derived from the following parts of the supersymmetric Lagrangian:

$$\mathcal{L}_Y = \lambda(S\psi_+\psi_- + \Phi_+\psi_S\psi_- + \Phi_-\psi_S\psi_+),$$
$$\mathcal{L}_g = -i\sqrt{2}g(\tilde{\Lambda}\psi_-\Phi_-^* + \tilde{\Lambda}\psi_+\Phi_+^*) + \mathrm{h.c.} \qquad (10)$$

We find that the gaugino, the inflatino and the higgsinos combine to form two Dirac spinors with masses given in Table IV. The one-loop effective potential reads:

$$V = \lambda^2 s^2|\Phi_-|^2 + \frac{g^2}{2}(|\Phi_-|^2 - \xi)^2 + \Delta V(s), \qquad (11)$$

Field	d.o.f.	squared mass
Φ_+	2	$\lambda^2 s^2 + g^2\xi$
Φ_-	2	$\lambda^2 s^2 - g^2\xi$
Dirac fermion	4	$\lambda^2 s^2$

TABLE III. Particle spectrum during D-term inflation, when $s \geq s_c$.

Field	d.o.f.	squared mass
S	2	$\lambda^2\hat{\phi}^2$
Φ_+	2	$\lambda^2\hat{\phi}^2 + 2\lambda^2 s^2$
$\sqrt{2}\Re(\phi_1)$	1	$2g^2\hat{\phi}^2$
A	3	$2g^2\hat{\phi}^2$
Dirac fermion	4	$(\frac{\lambda^2}{2}+g^2)\hat{\phi}^2 + \frac{\lambda^2}{2}s^2 + \sqrt{\Delta}$
Dirac fermion	4	$(\frac{\lambda^2}{2}+g^2)\hat{\phi}^2 + \frac{\lambda^2}{2}s^2 - \sqrt{\Delta}$

TABLE IV. Particle spectrum along the inflaton direction at the end of D-term inflation, with $s \leq s_c$ and $\hat{\phi} = (\xi - (\lambda/g)^2 s^2)^{1/2}$. For the fermion masses, we defined $\Delta \equiv (\frac{1}{2}\lambda^2(\hat{\phi}^2 + S^2) + g^2\hat{\phi}^2)^2 - 2\lambda^2 g^2\hat{\phi}^4$.

where $\Delta V(s)$ is given by Eq. (1). The supertrace vanishes at any time. The potential is continuous at $s = s_c$, and so is its derivative if we take:

$$\Lambda^2 = e^{\epsilon}\lambda^2 s_c^2, \qquad \epsilon \equiv \frac{1}{2} + \frac{\ln 2}{3}(1 - \frac{\lambda^2}{g^2}). \tag{12}$$

Then, for any choice of λ and g, the corrections are small with respect to the tree-level potential, and have the shape given in Fig. 1. Let us now comment on

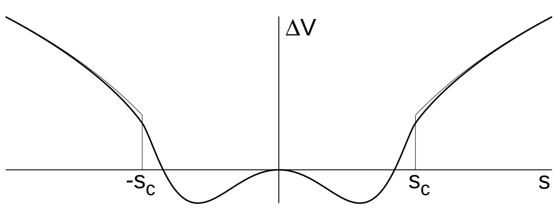

FIG. 1. Shape of the one-loop corrections, for F-term inflation with $g \leq \alpha$, and D-term inflation. The thin line shows the approximation of Eqs. (3,9) for $|S| \geq s_c$.

the importance of the one-loop corrections to the first and second derivative at the end of inflation. Obviously, they are dominant at the very begining of the symmetry breaking. At this time, it is important to know exactly $\partial\Delta V/\partial s$, in order to characterize the emergence of an effective Higgs background (from the coarse-graining of large-scale quantum fluctuations). After this stage, when the Higgs start(s) to grow, the loop corrections will affect $\partial V/\partial s$ and $\partial^2 V/\partial s^2$ only by a few percent[‡‡].

IV. CONCLUSION

In this paper, we have discussed the problem of supersymmetry breaking during and at the end of supersymmetric hybrid inflation, when the inflation scale is generated by GUT physics. We did not consider inflation models at intermediate or low energy scales [19,20]. Neither did we include supergravity corrections. When global supersymmetry is replaced by supergravity, the F-flat directions can only be preserved for specific Kähler potentials [3,6]; this does not apply to the case of D-term inflation, for which supergravity corrections are small for all values of the fields below the Planck mass [3,6]. In this framework, we calculated the one-loop corrections which modify the effective potential in the flat (inflaton) direction. We would like to point out that the classical trajectories of the Higgs fields do not necessarily coincide with their valley of local minima. However, for a wide range of parameters, the trajectory remains very close to this valley. Also, in more realistic cases, the Higgs fields will belong to non-trivial representations of G and their VEVs will break a non-abelian gauge group [22–25]. They may also couple to fermions such as right-handed neutrinos [12,23]. Hence, generally, one would find a much richer spectrum which would increase the corrections in a model-dependent way. Our results do not apply only to the end of inflation. They could be used in other models for which inflation takes place in local minima that spontaneously break both the gauge symmetry and supersymmetry. Also, in supersymmetric multiple inflationary models, part of the gauge symmetry breaks between the different stages of inflation, and at each step, the calculation of loop corrections can be performed as discussed in this paper.

[‡‡]For instance, in the D-term case, when the inflaton stabilizes around zero, its oscillation frequency is usually calculated from the tree-level effective mass: $\frac{\partial^2 V}{\partial s^2} = 2\lambda^2|\Phi_-|^2 = 2\lambda^2\xi$. As can be seen form Fig.1., the loop corrections will lower this value. We find:

$$\frac{\partial^2\Delta V}{\partial s^2} = -\frac{g^2\lambda^2\xi}{\pi^2} f(\frac{\lambda^2}{2g^2}),$$
$$f(x) \equiv \frac{\ln 2}{3}(1+x) - \frac{x\ln x}{2(1-x)}, \qquad f(1) \simeq 1.$$

So, with $\lambda = g = 1$, the effective mass is lowered by 4 %. A similar order of magnitude is found in the case of F-term inflation, for which $\frac{\partial^2 V}{\partial s^2} = 4\alpha\mu^2$:

$$\frac{\partial^2\Delta V}{\partial s^2} = -\frac{g^2\alpha\mu^2}{\pi^2}\tilde{f}(\frac{\alpha^2}{g^2}),$$

and $\tilde{f}(x)$ is a complicated function, of order one when $1 < x < 100$. When $\alpha = 2g = 1$, the effective mass is lowered by 2 %.

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